

x_F -dependence of J/Ψ suppression in pA collisions.

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Abstract

Coherence effects are important in the description of collisions with extended objects as nuclei. At fixed target energies and small x_F , the coherence length of the fluctuation containing the $c\bar{c}$ is small and the usual nuclear absorption model is valid. However, at higher energies and/or x_F the nucleus is seen as a whole by the fluctuation. In this case, the total, not the absorptive, $c\bar{c} - N$ cross section controls the suppression and also shadowing of gluons appears. We propose that the growth of the coherence length can explain the x_F -dependence of present experimental data. For this, we need a ratio of absorptive over total $c\bar{c} - N$ cross section of 0.2.

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The J/Ψ suppression is one of the main signals for quark gluon plasma (QGP) formation [1]. The *anomalous suppression* observed by NA50 Collaboration [2] is interpreted as produced by a deconfined state [3] though hadronic interpretations are possible [4]. Whether NA50 data give or not a definite proof of the formation of a deconfined state is a topic of intense discussion. In this letter we study the *non-anomalous suppression* in the whole range of x_F a subject that, though at first sight may seem solved, has some remaining open problems. This *normal suppression* is usually ascribed to multiple interaction of the produced pre-resonant $c\bar{c}$ state (color octet [5]) with the surrounding nuclear matter. The picture is very simple and well known: a $c\bar{c}$ is created at some point z_0 inside a nucleus in an octet state. In its travel through the nucleus, this state can interact at points $z > z_0$ with other nucleons that will destroy it with a cross section σ^{abs} . The formula describing this nuclear absorption is, after integrating in z :

$$\sigma_{pA} = \frac{\sigma_{pp}}{\sigma^{abs}} \int d^2b \left[1 - \exp \left(-\sigma^{abs} AT_A(b) \right) \right] \quad . \quad (1)$$

where $T_A(b) = \int dz \rho_A(b, z)$ is the profile function for nucleus A normalized to 1 and $\rho_A(b, z)$ is the nuclear density, that we take from ref. [6].

This formula describes well the observed pA data at midrapidities measured by NA38-NA51 [7] both for J/Ψ and Ψ' suppression, supporting the interpretation of the preresonant color octet state. A cross section of $\sigma^{abs} = 6.5 \pm 1.0$ mb is obtained by the experimental collaboration. This cross section does not depend on the energy. For this analysis, data from OCu , OU and SU collisions were also included. The E866/NuSea Collaboration [8] has also measured J/Ψ suppression in pA collisions and the result does not fully agree with the one from NA38-NA51: using the parametrization $\sigma^{pA} = A^\alpha \sigma^{pp}$ one obtains $\alpha \sim 0.95$ for E866/NuSea [8] at $y \sim 0$

and $\alpha \sim 0.92$ for CERN data. To reproduce E866 data at $x_F \sim 0$ a $\sigma^{abs} \sim 3$ mb is needed. E866/NuSea data give also the x_F dependence of the nuclear suppression, the origin of which can not be completely attributed to modifications of the nuclear gluon distributions [9] (see below) and remains an open problem. The main goal of the present work is to describe this x_F dependence and also to compare the different experiments. Understanding the rapidity pattern of the absorption is very important in order to have a real knowledge of the physics behind the suppression and also to extrapolate to RHIC and LHC energies, where this nuclear absorption would be present.

The idea is simple and its theoretical formulation has been derived in a previous paper [10]. In the frame where the nucleus is at rest, the incoming proton fluctuates in a complicated system of quarks and gluons with coherence length l_c . At small energies and x_F , l_c is small (of the order of the nucleon size) and only one nucleon in the nucleus takes part in the hard interaction that produces the $c\bar{c}$. This implies $\sigma_{pA} \sim A$. This behavior is modified by the collisions of the produced $c\bar{c}$ with the other nucleons in the nucleus. We have, in this way, the usual description of nuclear absorption given by eq. (1). However, at large energies and/or x_F , l_c gets eventually larger than the nuclear size and the nucleus is seen as a whole by the fluctuation. As a consequence the time ordering is lost and (1) is no longer valid. To describe this regime, we have introduced two types of collisions with the nucleons in the target, the ones of the light partons (mainly gluons) – with total cross section σ – and the ones of the heavy system ($c\bar{c}$) – with total cross section $\tilde{\sigma}$ – in an eikonal approach. The first ones give rise to modifications in the nuclear gluon distribution and the second to a suppression of charmonia states (the generalization of nuclear absorption). The result which replaces eq. (1) is [10]

$$\frac{d\sigma_{pA}}{dx_F d^2b} = \sigma_{pQCD}^{gg \rightarrow c\bar{c}X} g_p(x_1, Q^2) g_A(x_2, Q^2, b) e^{-\frac{1}{2}\tilde{\sigma} AT_A(b)}, \quad (2)$$

with $x_F = x_1 - x_2$, $x_1 x_2 s = m_{J/\Psi}^2$, and

$$g_A(x_2, Q^2, b) = 2 \int d\omega \left[1 - e^{-\frac{1}{2}\sigma(\omega) AT_A(b)} \right]. \quad (3)$$

ω represents kinematical variables of the gluon-nucleon interactions to be integrated. This term gives the modification of nuclear structure functions with respect to nucleons and will not be discussed here – see for instance [11] for a model. In (2) we see that the multiple scatterings of the heavy and light systems factorize, so that we can separate both contributions in the ratios R_{pA} of pA to pp cross sections.

$$R_{pA} = R_{pA}^{shadow} R_{pA}^{c\bar{c}}. \quad (4)$$

Neglecting shadowing corrections to gluons, the change from low to asymptotic energies consists in the substitution:

$$\frac{1}{\sigma^{abs}} \left[1 - \exp \left(-\sigma^{abs} A T_A(b) \right) \right] \longrightarrow A T_A(b) \exp \left(-\frac{1}{2}\tilde{\sigma} A T_A(b) \right), \quad (5)$$

The most important point is the change of the absorptive cross section by the total one. When $\tilde{\sigma} = \sigma^{abs}$ the first correction term in the expansion in $\tilde{\sigma}$ is the same for both expressions. As these cross sections are not very large in practice, the numerical values turn out to be very similar. In this case, this justifies the use of formula (1) for high energy though strictly speaking it is only valid at small energies. In our case, however, we will suppose that $\tilde{\sigma} \neq \sigma^{abs}$. That is, the preresonant $c\bar{c}$ state has a non negligible probability of not being destroyed. This increases the absorption for large values of x_F where coherence is reached and $\tilde{\sigma}$ has to be used.

The two regimes are particular solutions of a more general equation which takes into account the coherence effects for any l_c . In this case the factorization given in (4) is no longer valid. Neglecting shadowing to structure functions,

$$\frac{d\sigma_A}{dx_F} = \sigma_{pQCD}^{gg \rightarrow c\bar{c}X} g_p(x_1, Q^2) g_p(x_2, Q^2) \sum_{n=1}^A C_A^n \sum_{j=1}^n \int d^2b T_n^{(j)}(b) \sigma_n^{(j)} \quad , \quad (6)$$

where

$$\sigma_n^{(j)} = j \left(-\frac{\tilde{\sigma}}{2} \right)^{j-1} (-\sigma^{abs})^{n-j} + (j-1) \left(-\frac{\tilde{\sigma}}{2} \right)^{j-2} (-\sigma^{abs})^{n-j+1} \quad . \quad (7)$$

The powers of nuclear profile functions in the expansion of (2) have to be changed to

$$T^n(b) \longrightarrow T_n^{(j)}(b) = n! \int_{-\infty}^{+\infty} dz_1 \int_{z_1}^{+\infty} dz_2 \dots \int_{z_{n-1}}^{+\infty} dz_n \cos(\Delta(z_1 - z_j)) \prod_{i=1}^n \rho_A(b, z_i). \quad (8)$$

The physical interpretation of the above equations is very clear, the first $j-1$ collisions are coherent (one of them corresponding to the light amplitude σ), the last $n-j$ are not and the j -th has the two possibilities. The Δ factor that controls these finite energy effects can be approximated by

$$\Delta \equiv \frac{1}{l_c} = \frac{m_p M^2}{s x_1} \quad , \quad (9)$$

where m_p is the mass of the proton and M^2 is the effective mass of the fluctuation. This effective mass could be measured in diffractive events containing a J/Ψ . In order to obtain the observed scaling of the suppression in x_F , we take $M^2/s \sim$ constant. This means that increasing the energy allows the fluctuation to have bigger masses. This is the kind of behavior of triple pomeron contribution to diffraction [12], however it is not clear that these are the relevant diagrams in the present case.

We take this scaling as a phenomenological ansatz, having in mind that this is most probably a finite energy effect that would disappear at asymptotic energies[‡]. In practice, what we have done is taking $M^2 = m_{J/\Psi}^2$ for the smallest experimental energy ($\sqrt{s} \sim 20$ GeV) and use the proportionality with s for the others.

In our calculations we have also introduced the nuclear corrections to parton distributions (gluons) as given by the EKS98 parametrization [9]. We could also use eq. (3), however this would need some theoretical inputs from gluon-nucleon cross section that complicates the computation. Moreover, gluon antishadowing can not be reproduced by this formula. To take into account coherence effects we have defined

$$R_{pA}^{shadow} = \left[1 - \left(1 - \frac{A_{eff}^{fin} - A_{eff}^{asym}}{A_{eff}^{prob} - A_{eff}^{asym}} \right) (1 - R_{pA}^{EKS}) \right] \quad (10)$$

R_{pA}^{EKS} being the shadowing corrections to $c\bar{c}$ production computed with EKS98 parametrizations, and A_{eff}^{prob} , A_{eff}^{asym} and A_{eff}^{fin} the effective A (i.e. σ_{pA}/σ_{pp}) computed with eq. (1), (2) and (6) respectively, all without shadowing. In this way, $R_{pA}^{shadow} = 1$ when l_c is very small ($A_{eff}^{fin} = A_{eff}^{prob}$) and $R_{pA}^{shadow} = R_{pA}^{EKS}$ for large coherence lengths ($A_{eff}^{fin} = A_{eff}^{asym}$). However, as the data from E866/NuSea uses Be as a reference, the shadowing corrections are not large and, in any case, much smaller than the observed effect (see Fig. 1). The final result for absorption is given by eq. (4) with $R_{pA}^{c\bar{c}} = A_{eff}^{fin}/A$.

With all these ingredients we have fitted the data from E866/NuSea Collaboration. The free parameters are σ^{abs} and $\tilde{\sigma}$. To have a good description of E866/NuSea data we obtained $\sigma^{abs} = 3$ mb and $\tilde{\sigma} = 15$ mb. The comparison with experimental

[‡]I thank A. Capella and A. Kaidalov for discussions on this point.

data can be seen in Fig. 1[§]. Also shown in this figure is the comparison taking into account only shadowing corrections to gluons given by EKS98 parametrization and nuclear suppression given by eq. (1) with $\sigma^{abs}=4.5$ mb. Notice that if shadowing (antishadowing at $x_F \sim 0$) is not included, $\sigma^{abs} \sim 3$ mb is needed in order to reproduce the data at $x_F \lesssim 0.2$.

In Fig. 2 we present the comparison with NA38-NA51 data. These data are measured in the interval $3 < y_{lab} < 4$. We have used the values $\langle x_F \rangle = 0.03$ for $E_{lab} = 450$ GeV data and $\langle x_F \rangle = 0.16$ for $E_{lab} = 200$ GeV data. The description of the data is good in spite of the apparent discrepancy in the parameter α among the two sets discussed above. Notice that the data at 200 GeV is more suppressed due to the larger x_F .

Finally, in Fig. 3 comparison is made with NA3 data [13] at $E_{lab} = 200$ GeV. The description is again not bad, though some discrepancy appears at large values of x_F . This is also observed in other analysis [14] and could be a signal of energy loss [15], however, the evidence is too weak and more experimental data would be needed in this region[¶]. Besides, including shadowing corrections (in fact antishadowing in this region) makes our results to increase, breaking the scaling on x_F mentioned above. This corrections are very uncertain and this could also be the origin of the discrepancy.

Let's compare our analysis with previous ones. In [14] a description of data very similar to ours was obtained by taking into account two different cross sections for octet and singlet $c\bar{c}$ states, the drawback of this model was the smallness of the lifetime of the octet state. In [17], some kind of energy loss of gluons from

[§]Actually, only W/Be data have been used in the fit, the Fe/Be data came out as a result.

[¶]Notice that energy loss with nuclear matter becomes smaller with increasing energy [16], so the effect in E866/NuSea will be smaller.

the incoming proton is proposed. In [18] different effects as shadowing, energy loss, comover absorption and intrinsic charm are taken into account separately. In [19] a model similar to ours, that takes into account coherence effects, has been developed. They also include two types of interactions for light and heavy systems, obtaining shadowing and $c\bar{c}$ suppression, however, their x_F dependence come from loosing of momentum of the $c\bar{c}$ pair as it experiences multiple scattering. Finally, in [20] coherence effects as well as other effects as energy loss and time formation are taken into account. This is a more formal analysis, but where coherence give effects similar to ours. Their computations are only for the χ_c state. One main difference is that they don't take into account an effective mass of the fluctuation depending on the energy. This makes their comparison with E866 data less good than ours (we will obtain a similar result by fixing $M^2 = M_{J/\Psi}^2$). In a previous work [21] their comparison with E866 data where better but they limited their analysis to $-0.1 < x_F < 0.25$ in order to see the influence of the formation time.

To conclude, a high energy pA collision can be seen, in the laboratory frame, as the multiple scattering of the system of quarks and gluons into which the incoming proton fluctuates. The coherence length of this fluctuation increases with the energy and x_F and two regions can be distinguished. When $l_c \ll R_A$ no shadowing to gluons is present and the suppression of $c\bar{c}$ pairs is given by the usual formula (1). When $l_c \gtrsim R_A$ the whole nucleus takes part in the collision and typical coherent phenomena as shadowing appear. The main point is that this change of regime is accompanied by a change $\sigma^{abs} \rightarrow \tilde{\sigma}$ in the expressions. If these two cross sections are different, the suppression at large values of x_F and/or larger energies is bigger. Let's also comment about the Ψ' : experimental data from NA38 don't see any difference between J/Ψ and Ψ' suppression, whereas some difference seems to appear in E866/NuSea. This

difference would be very easily accounted for in our approach just by taking the absorptive cross sections different for singlet and octet $c\bar{c}$ states: as the total J/Ψ has a contribution of $\sim 40\%$ singlet coming from disintegrations of χ_c , a larger σ^{abs} for octet than for singlet would explain the difference. This would introduce a new parameter. Finally, the extrapolations to RHIC and LHC energies depend on two assumptions, the energy dependence of $\tilde{\sigma}$ and σ^{abs} (that we have taken as constants) and the value M^2 of the effective mass of the fluctuation. With this last assumption we are wondering if the x_F scaling of the suppression will still be valid at high energy (in the case $M^2/s \sim \text{const.}$) or if, on the contrary, a scaling in x_2 will appear, this last possibility is the most reasonable as the proportionality in s of M^2 seems to be a finite energy effect. Let's give some estimations for RHIC and LHC at $x_F \sim 0$. Assuming no energy dependence for the cross sections $\tilde{\sigma}$ and σ^{abs} we obtain for pAu collisions at $\sqrt{s}=200$ GeV a ratio over pp of 0.81 if the scaling in x_F is maintained and 0.43 if it is not. This means ratios of 0.67 and 0.17 for $AuAu$ collisions respectively. In the case of LHC, the only difference is shadowing, that will not affect the case $M^2/s \sim \text{const.}$, in the other case, we obtain 0.31 for pPb collisions and 0.1 for $PbPb$ at $\sqrt{s}=5500$ GeV.

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Figure captions:

Fig. 1. Comparison of our results on x_F dependence of nuclear suppression (solid lines) with E866 data [8] for Fe/Be and W/Be . Effects of shadowing as given by EKS98 parametrization including absorption with $\sigma^{abs} = 4.5$ mb in eq. (1) are also shown (dashed lines).

Fig. 2. $B_{\mu\mu}\sigma^\psi/A$ measured by NA38 Collaboration [7] at $E_{lab}=450$ GeV² (black boxes) and $E_{lab}=200$ GeV² (white boxes) compared with our results at $E_{lab}=450$ GeV² (solid line) and $E_{lab}=200$ GeV² (dotted line). Also shown, the suppression given by eq. (1) with $\sigma^{abs}=3$ mb.

Fig. 3. α parameter as a function of x_F compared with data from NA3 [13] at $E_{lab}=200$ GeV².

Figure 1

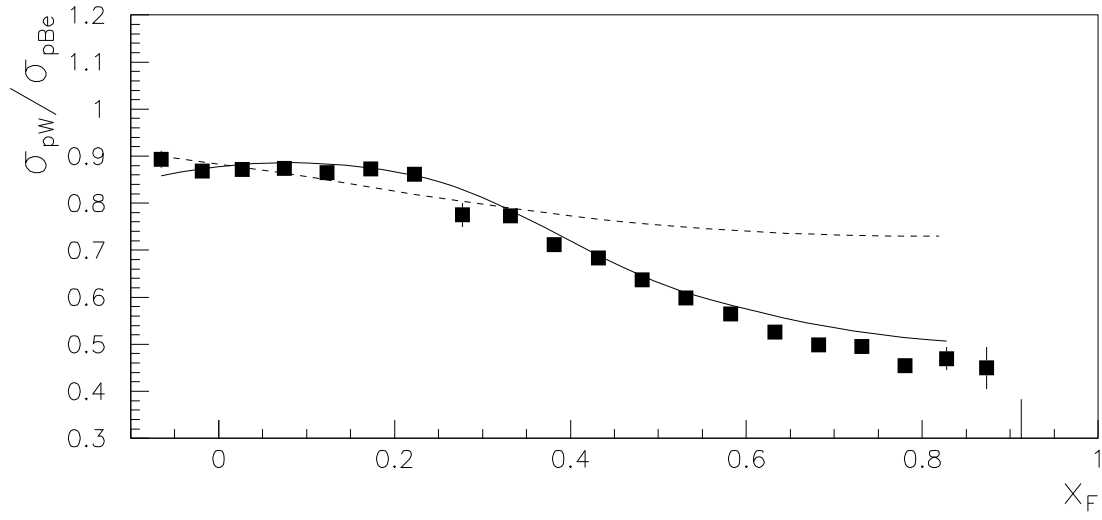
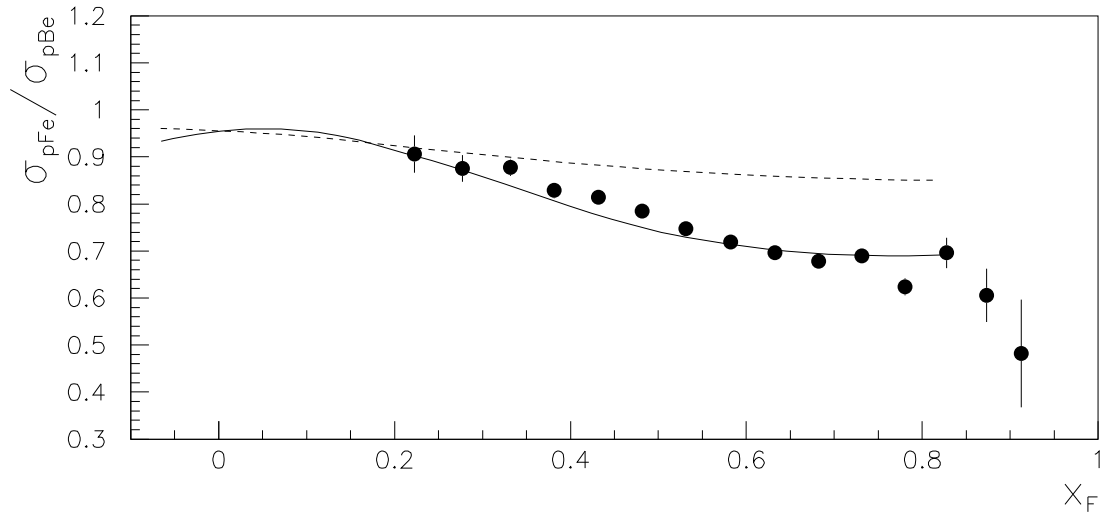


Figure 2

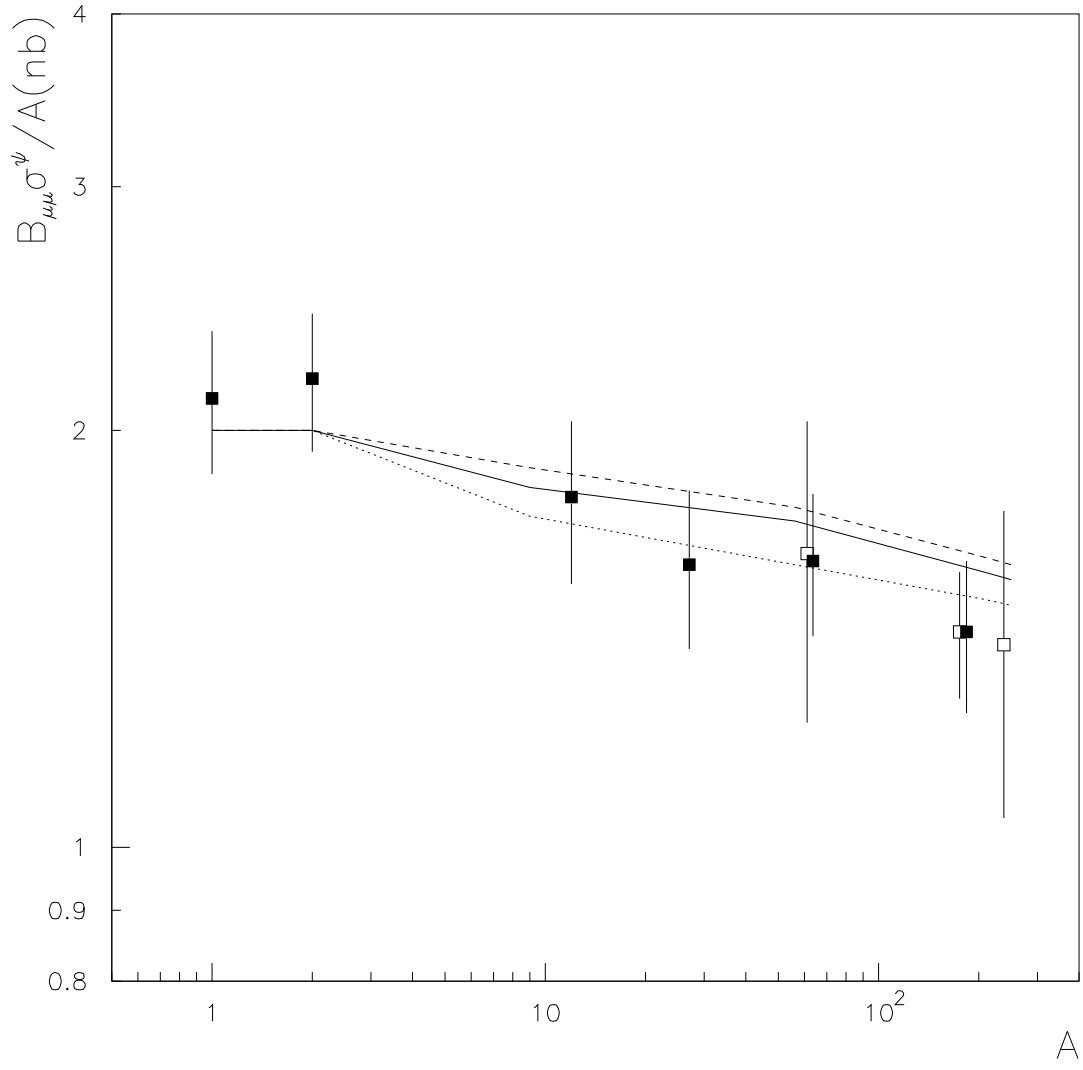


Figure 3

